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Baryonic States in QCD From Gauge/String Duality at Large N_C ^{*}

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Abstract

We have computed the baryon spectrum in the context of $\mathcal{N} = 4$ superconformal Yang-Mills theory using AdS/CFT duality. Baryons are included in the theory by adding an open string sector, corresponding to quarks in the fundamental and higher representations. The hadron mass scale is introduced by imposing boundary conditions at the wall at the end of AdS space. The quantum numbers of each baryon, are identified by matching the fall-off of the string wavefunction $\Psi(x, r)$ at the asymptotic 3+1 boundary to the operator dimension of the lowest three-quark Fock state, subject to appropriate boundary conditions. Higher Fock states are matched quanta to quanta with quantum fluctuations of the bulk geometry about the fixed AdS background, maintaining conformal invariance. The resulting four-dimensional spectrum displays a remarkable resemblance to the physical baryon spectrum of QCD, including the suppression of spin-orbit interactions.

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1 Introduction

An outstanding consequence of Maldacena's duality [1] between 10-dimensional string theory on $AdS_5 \times S^5$ and Yang-Mills theories at its conformal 3+1 spacetime boundary [2, 3] is the potential to describe processes for physical QCD which are valid at strong coupling and do not rely on perturbation theory. As shown by Polchinski and Strassler [4], dimensional counting rules [5] for the leading power-law fall-off of hard exclusive glueball scattering can be derived from a gauge theory with a mass gap dual to supergravity in warped spacetimes. The modified theory generates the hard behavior expected from QCD, instead of the soft behavior characteristic of strings. Other examples are the description of form factors at large transverse momentum [6] and deep inelastic scattering [7]. The discussion of scaling laws in warped backgrounds has also been addressed in [8, 9, 10].

The AdS/CFT duality gives a non-perturbative definition of quantum gravity in a curved background which is asymptotic to a product of Anti-de Sitter space AdS and a compact Einstein space X . As originally formulated [1], a correspondence was established between the supergravity approximation to Type IIB string theory, and the large N_C brane decoupling limit with gauge dynamics corresponding to $\mathcal{N} = 4$ super Yang-Mills (SYM) in four dimensions. The bulk geometry has exact conformal geometry $AdS_5 \times S^5$ in the near-horizon region $r \ll R$, where $R = (4\pi g_s N_C)^{1/4} \alpha'^{1/2}$ is the radius of AdS and the radius of the five-sphere. The extra five dimensions of S^5 correspond to the $SU(4)$ global symmetry which rotates the particles present in the SYM supermultiplet in the adjoint representation of $SU(N_C)$. The conformal group $SO(2, 4)$ is identified with the isometry group of AdS_5 , and $SU(4) \simeq SO(6)$ with the isometries of S^5 . The supergravity duality requires a large value of R corresponding to a large value of the 't Hooft parameter, $g_s N_C$ [11].

In a recent attempt to extend the glueball results to hadrons [12], we used the AdS/CFT correspondence to determine the basic properties of hadronic light-front wavefunctions (LFWF) in QCD based on the underlying conformal symmetry of the duality for pointlike hard-scattering processes which occur in the large- r region of AdS space. The scaling behavior of the string modes determine the behavior of the QCD hadronic wavefunction, giving a precise counting rule for each Fock component state with an arbitrary number of quarks and gluons and internal orbital angular momentum¹. The discussion was carried out in terms of the lowest dimensions of interpolating fields near the boundary of AdS, treating the boundary values of the string states $\Psi(x, r)$ as a product of quantized operators which create n -partonic states

¹In [12] we examined the possibility of identifying the internal orbital momentum of hadrons with Kaluza-Klein excitations of the internal space S^5 . Henceforth we follow the interpretation given here in terms of quantum fluctuations about the AdS background.

out of the vacuum [12]. Our AdS/CFT derivation validate QCD perturbative results and confirm the dominance of the quark interchange mechanism [13] for exclusive QCD processes at large N_C . The predicted orbital dependence coincides with the fall-off of light-front Fock wavefunctions derived in perturbative QCD [14]. Since all of the Fock states of the LFWF beyond the valence state are a manifestation of quantum fluctuations, it is natural to match quanta to quanta the additional dimensions with the metric fluctuations of the bulk geometry about the fixed AdS background.

For large values of R , or small curvature of AdS space, it is expected that the dual of a Yang Mills theory is classical gravity. The correspondence also implies that the dual of strongly coupled QCD is a weakly coupled string. Since QCD is weakly coupled at high energies, the dual theory is expected to be a strongly coupled string model at small 't Hooft coupling and would require the understanding of strings in highly curved backgrounds, extending the semiclassical approximation to include quantum effects on the string theory side. The behavior of string states in the infrared region is dependent on dynamics at small- r , and it is a priori unknown. Non-conformal aspects are needed to make contact with the real world. The introduction of quarks in the fundamental representation is also crucial, requiring an open string sector.

In spite of the difficulties mentioned above, important progress has been achieved by extending the AdS/CFT correspondence beyond the supergravity approximation to construct string duals to non-conformal gauge theories². Even if the detailed form of the metric at small- r is unknown, salient QCD dynamical features, such as the generation of a mass gap and a hadron spectrum, will follow from the deformation of the AdS conformal background at small r . Indeed from [4] there follows a simple relation between the 10-dimensional string scale α'_s and the Yang-Mills 4-dimensional scale α'_{QCD} in a warped space: $\alpha'_{QCD} \sim \alpha'_s (R/r_o)^2$, where the cutoff $r_o = \Lambda_{QCD} R^2$, breaks conformal invariance and allows the introduction of the QCD scale.

A physical hadron in four-dimensional Minkowski space has four-momentum P_μ and hadronic invariant mass states given by $P_\mu P^\mu = \mathcal{M}^2$. The string wavefunction in r is the extension of the baryon wavefunction into the fifth dimension: we thus analytically match the three-quark proton wavefunction in $3+1$ space to its corresponding string wavefunction using the 10-dimensional wave equation. Different values of r correspond to different energy scales at which the hadron is examined. In particular, the $r \rightarrow \infty$ boundary corresponds to the $Q \rightarrow \infty$, zero separation limit. The physical string modes

$$\hat{\Psi} \equiv \langle 0 | \hat{\Psi} | P \rangle \sim e^{-iP \cdot x} f(r) Y(y), \quad (1)$$

are plane waves along the Poincaré coordinates, and Y is a function of the transverse coordinates y . For large- r , $f(r)$ scale as $f(r) \sim r^{-\Delta}$, where Δ is the conformal di-

²For a review see [15]. See also [11].

mension of the string state, the same dimension of the interpolating operator which creates a hadron out of the vacuum [4]. The string modes are coupled to the matter fields of the conformal theory as determined by the boundary limit between the string partition function and the generating functional of the quantum field theory [2, 3]. For example, the quantum numbers of each baryon, including intrinsic spin and orbital angular momentum, are determined by matching the dimensions of the string modes $\Psi(x, r)$, with the lowest dimension of the baryonic interpolating operators in the conformal limit. Although the underlying string theory dual to $QCD(3 + 1)$ is unknown, the introduction of quark fields in the fundamental representation and the symmetries in the asymptotic boundary, should provide the conditions required to establish a precise matching between the string modes in the semiclassical gravity approximation and boundary states with well defined number of partons. In particular, the known light baryons are dual to spin- $\frac{1}{2}$ and spin- $\frac{3}{2}$ strings and consequently there is not a vastly large mass gap between baryons with total angular momentum $J \leq 2$ and $J > 2$ for large N_C .

After stating in Sec. 2 some basic properties of the correspondence between string modes in AdS space and baryon states at the asymptotic boundary, we confront our results with the spectrum of nucleon and Δ orbital resonances in Sec. 3. Some concluding remarks are given in Sec. 4.

2 Baryon Interpolating Operators and String Modes

A precise statement of the duality between a string/gravity theory on a $(d + 1)$ -dimensional Anti-de Sitter space AdS_{d+1} and the large N_C limit of a conformal theory at its d -dimensional boundary, is given formally in terms of the full partition function of the string theory in the bulk Z_{string} which should coincide with the generating functional of the conformal field theory Z_{CFT} on the AdS boundary [2, 3]: $Z_{string}[\Psi(x, z = 0)] = Z_{CFT}[\Psi_o]$. For spin- $\frac{1}{2}$ dilatino in the bulk, the duality involves positive and negative chirality components [16] $\Psi^\pm = \frac{1}{2}(1 \pm \gamma_5)\Psi$, which couple with CFT operators \mathcal{O}^+ and \mathcal{O}^-

$$Z_{CFT}[\Psi_o, \bar{\Psi}_o] = \left\langle \exp \left(i \int d^d x \left[\bar{\mathcal{O}}^- \Psi_o^- + \bar{\Psi}_o^+ \mathcal{O}^+ \right] \right) \right\rangle. \quad (2)$$

Near the boundary of AdS, $z = R^2/r \rightarrow 0$, the independent solutions of the 10-dimensional Dirac equation are

$$\begin{aligned} \Psi(z, x) &\rightarrow z^{\frac{d}{2}+mR} \Psi_+(x) + z^{\frac{d}{2}-mR} \Psi_-(x) \\ \bar{\Psi}(z, x) &\rightarrow z^{\frac{d}{2}+mR} \bar{\Psi}_+(x) + z^{\frac{d}{2}-mR} \bar{\Psi}_-(x). \end{aligned}$$

The solution with Ψ_- dominates near $z \rightarrow 0$, thus $\Psi_- = \Psi_o$. The field $\Psi_-(x)$ acts as a boundary source, and $\Psi_+(x)$ is the response function which incorporates the quantum fluctuations. The boundary sources for positive and negative chirality Ψ_o^- and $\overline{\Psi}_o^+$ have dimensions $\frac{d}{2} - mR$. Consequently the dimension of the CFT operators \mathcal{O}^+ and $\overline{\mathcal{O}}^-$ is $\frac{d}{2} + mR$. Since the dimension of Ψ_+ is also $\frac{d}{2} + mR$, we expect that $\Psi_+(x)$ is related to the expectation value of \mathcal{O} in the presence of the source Ψ_o : $\Psi_+(x) \sim \langle \mathcal{O} \rangle_{\Psi_o}$. Indeed $\langle \mathcal{O} \rangle = (2\Delta - d) \Psi_+(x)$ [17]. Thus Ψ_+ acts as a semiclassical field and is the boundary limit of the normalizable string solution.

We consider first the classical solution dual to the valence Fock state, described by the massless 10-dimensional Dirac equation in the bulk: $\Gamma^A D_A \hat{\Psi} = 0$. Full space coordinates are $x^A = (x^\mu, z, y^a)$, with x^μ the Minkowski variables and z the holographic coordinate. Coordinates in the compact space are y^a , and $g_{\perp ab}$ is the transverse metric. The full metric of spacetime is [4]

$$ds^2 = \frac{R^2}{z^2} e^{2A(z)} (\eta_{\mu\nu} dx^\mu dx^\nu - dz^2) + g_{\perp ab} dy^a dy^b, \quad (3)$$

where $A(z) \rightarrow 0$ as $z \rightarrow 0$, and behaves asymptotically as a product of AdS space and a compact manifold X . A 10-dimensional field is represented by hat quantities: $\hat{\Phi}, \hat{\Psi}$; a field on AdS space by Φ, Ψ , and ϕ, ψ represent fields in 4-dimensions. Knowledge of the full geometry is required to solve the Dirac equation. We truncate effectively the infrared region at the end of AdS space at $z_o = 1/\Lambda_{QCD}$, where string modes cannot propagate. We expand the state $\hat{\Psi}$ in terms of spinors $\eta(y)$ of the Dirac operator on a $d+1$ sphere with eigenvalues λ_κ as $\hat{\Psi}(x, z, y) = \sum_\kappa \Psi_\kappa(x, z) \eta_\kappa(y)$. For each eigenvalue λ the normalizable string modes are

$$\Psi(x, z) = C e^{-iP \cdot x} z^{\frac{d+1}{2}} \left[J_{\lambda R - \frac{1}{2}}(z\mathcal{M}) \mu_+(P) + J_{\lambda R + \frac{1}{2}}(z\mathcal{M}) \mu_-(P) \right]. \quad (4)$$

Four dimensional spinors are related by [16] $\mu_- = (\not{P}/P) \mu_+$. The eigenvalues on S^{d+1} are $\lambda_\kappa R = \pm (\kappa + \frac{d}{2} + \frac{1}{2})$, $\kappa = 0, 1, 2, \dots$, with multiplicity [18] $D_{d+1}(\kappa) = 2^{\frac{d}{2}} \binom{d+\kappa}{\kappa}$.

The lowest dimension of a spin- $\frac{1}{2}$ field on AdS_{d+1} space is $\text{Dim}[\Psi] = d + \frac{1}{2}$. A spin- $\frac{1}{2}$ field in the d -dimensional boundary theory has dimensions $\text{Dim}[\psi] = (d-1)/2$. It is quite remarkable that a dual CFT boundary operator \mathcal{O} with dimension $d + \frac{1}{2}$ can only be constructed for $d = 2, 4$ and $d \rightarrow \infty$ with a product of five, three and two quark fields respectively, to match the string dimensions in the bulk.

For $d = 4$, the spinor irreps are **4**, **20**, **60**, **140**, \dots Classical spin- $\frac{1}{2}$ string solutions are labeled by eigenvalues of the Dirac operator on X . The lowest string mode for $\kappa = 0$ has dimension $\Delta = \frac{9}{2}$, and transforms as a **4** under the $SU(4)$ R -symmetry. The corresponding CFT operator $\mathcal{O}_{\frac{9}{2}}$ is constructed as the product of three quark fields ψ^r transforming as a **4** of $SU(4)$, since $\overline{\mathbf{4}} \otimes \overline{\mathbf{4}} \otimes \overline{\mathbf{4}} \rightarrow \mathbf{4}$. The new degeneracy

could be interpreted as a flavor symmetry with $r = u, d, c, s$, which is broken by quark masses. With respect to color $\mathcal{O}_{9/2}$ has the gauge invariant form $\mathcal{O}(x)_{9/2} = \psi_{\mathbf{N}_C}(x)\psi_{\mathbf{N}_C}(x)\psi_{\mathbf{N}_C(\mathbf{N}_C-1)/2}(x)$, where the representation $\mathbf{N}_C(\mathbf{N}_C-1)/2$ follows from the antisymmetric component of the tensor product $\mathbf{N}_C \otimes \mathbf{N}_C$. For $N_C = 3$, we recover the usual form of the interpolating operator which creates a physical baryon in QCD(3+1): $\mathcal{O}_{9/2} = \epsilon_{abc}\psi_a\psi_b\psi_c$.

QCD is fundamentally different from SYM theories where matter fields appear in adjoint multiplets of $SU(N_C)$. The $\mathcal{N} = 4$ theory is dual to the low energy supergravity approximation to type IIB string [1] compactified on $AdS_5 \times S^5$. The SYM fields correspond to closed strings and comprise a gluon field, six scalars, four Majorana gluinos, and their antiparticles. The introduction of quarks in the fundamental representation³ is dual to the introduction of an open string sector, where the strings end on a brane and join together at a point in the AdS geometry [19]. The SYM particles are expected to acquire a mass of the order of SUSY breaking scale and decouple from the theory.

The AdS/CFT correspondence is interpreted in the present context as a classical duality between the lowest three-quark valence state in the asymptotic $3 + 1$ boundary and the lowest string mode in $AdS_5 \times S^5$. Higher Fock components are a manifestation of the quantum fluctuations of QCD and are conformal states in the limit of massless quarks and vanishing QCD β -function [22]. Metric fluctuations of the bulk geometry about the fixed AdS background should correspond to quantum fluctuations of Fock states above the valence state in the limit where QCD appears nearly conformal. As shown by Gubser, Klebanov and Polyakov, orbital excitations in the boundary correspond to string degrees of freedom propagating in the bulk from quantum fluctuations in the AdS sector [23].

Consider a gauge invariant interpolating operator which creates an arbitrary Fock-state in the boundary

$$\mathcal{O}_n^L = \text{Tr} (\psi D_{\ell_1} \psi D_{\ell_2} \psi \dots \bar{\psi} \psi \dots F D_{\ell_m} F \dots), \quad (5)$$

where the m derivatives determine the total spacetime orbital angular momentum, $L = \sum_{i=1}^m \ell_i$. The conformal dimension of (5) is $\Delta = \Delta_n + L$, where Δ_n is the sum of the dimensions of the quarks, antiquarks and gluons of the Fock component. An effective five-dimensional mass μ in the AdS wave equation corresponding to quantum excitations about the fixed AdS metric is asymptotically determined by spectral comparison of the string modes and the boundary operators of QCD, while maintaining conformal invariance. Matching the dimension of the Fock components at

³The introduction of a finite number of N_f branes is dual to the introduction of flavor [20] with quarks in the fundamental representation, and leads to a calculable spectrum [21].

the Minkowski boundary we find the relation $\mu R = \hbar c L$. Thus, an ℓ quantum orbital at the Minkowski boundary corresponds to a five-dimensional mass $\mu \sim \hbar c \ell / R$ in the bulk side. The four-dimensional mass spectrum \mathcal{M}_L is determined by imposing boundary conditions on one of the solutions of the Dirac equation $\Psi^\pm(x, z_o) = 0$. The solution of the spin- $\frac{3}{2}$ Rarita-Schwinger equation in AdS space is more involved, but considerable simplification occurs for polarization along Minkowski coordinates, Ψ_μ , where it becomes similar to the spin- $\frac{1}{2}$ solution [24].

3 Baryon Spectrum

The study of the hadron spectrum is crucial for our understanding of bound states of strongly interacting relativistic confined particles. Different QCD-based models often disagree, even in the identification of the relevant degrees of freedom [25, 26]. Studies of orbitally excited baryons based on the $1/N_C$ expansion have been useful for identifying the relevant effective operators and determine their relative importance [27]. An outstanding puzzle is that the spin-orbit splitting, which experimentally is very small, appear in the $1/N_C$ expansion as a zeroth-order effect. Recently, the computation of orbital excitations on the lattice has been extended up to spin- $\frac{5}{2}$ states [28].

The spectrum of N and Δ baryon states is listed in Table 1 according to total angular momentum-parity assignment given by the PDG [29]. To determine intrinsic spin and orbital momentum quantum numbers we have used the conventional $SU(6) \supset SU(3)_{flavor} \otimes SU(2)_{spin}$ multiplet structure. We limit ourselves to the light unflavored hadron states and the introduction of massless quarks. Since $m_{u,d} \ll \Lambda_{QCD}$, the light quarks are extremely relativistic. Consequently the mass of the hadrons corresponds essentially to the confined kinetic energy of massless quarks and gluons. The intrinsic spin S of a given hadron should match the spin of the dual string.

We present in Fig. 1 the orbital spectrum of the nucleon states and in Fig. 2 the Δ orbital resonances. We plot the values of \mathcal{M}^2 as a function of L . The nucleon states with intrinsic spin $S = \frac{1}{2}$ lie on a curve below the nucleons with $S = \frac{3}{2}$. We have chosen our boundary conditions by imposing the condition $\Psi^+(x, z_o) = 0$ on the positive chirality modes for $S = \frac{1}{2}$ nucleons, and $\Psi_\mu^-(x, z_o) = 0$ on the chirality minus strings for $S = \frac{3}{2}$. In contrast to the nucleons, all of the known Δ orbital states with $S = \frac{1}{2}$ and $S = \frac{3}{2}$ lie on the same trajectory. The boundary conditions in this case are imposed on the chirality minus string modes. The numerical solution corresponding to the roots of Bessel functions in (4), give the nonlinear trajectories indicated in the figures. All the curves correspond to the value $\Lambda_{QCD} = 0.22$ GeV, which is the only

$SU(6)$	S	L	Baryon State			
56	$\frac{1}{2}$	0	$N_{\frac{1}{2}}^{1+}(939)$			
	$\frac{3}{2}$	0	$\Delta_{\frac{3}{2}}^{3+}(1232)$			
70	$\frac{1}{2}$	1	$N_{\frac{1}{2}}^{1-}(1535)$	$N_{\frac{3}{2}}^{3-}(1520)$		
	$\frac{3}{2}$	1	$N_{\frac{1}{2}}^{1-}(1650)$	$N_{\frac{3}{2}}^{3-}(1700)$	$N_{\frac{5}{2}}^{5-}(1675)$	
	$\frac{1}{2}$	1		$\Delta_{\frac{1}{2}}^{1-}(1620)$	$\Delta_{\frac{3}{2}}^{3-}(1700)$	
56	$\frac{1}{2}$	2		$N_{\frac{3}{2}}^{3+}(1720)$	$N_{\frac{5}{2}}^{5+}(1680)$	
	$\frac{3}{2}$	2	$\Delta_{\frac{1}{2}}^{1+}(1910)$	$\Delta_{\frac{3}{2}}^{3+}(1920)$	$\Delta_{\frac{5}{2}}^{5+}(1905)$	$\Delta_{\frac{7}{2}}^{7+}(1950)$
70	$\frac{1}{2}$	3		$N_{\frac{5}{2}}^{5-}$	$N_{\frac{7}{2}}^{7-}$	
	$\frac{3}{2}$	3	$N_{\frac{3}{2}}^{3-}$	$N_{\frac{5}{2}}^{5-}$	$N_{\frac{7}{2}}^{7-}(2190)$	$N_{\frac{9}{2}}^{9-}(2250)$
	$\frac{1}{2}$	3		$\Delta_{\frac{5}{2}}^{5-}$	$\Delta_{\frac{7}{2}}^{7-}$	
56	$\frac{1}{2}$	4		$N_{\frac{7}{2}}^{7+}$	$N_{\frac{9}{2}}^{9+}(2220)$	
	$\frac{3}{2}$	4	$\Delta_{\frac{5}{2}}^{5+}$	$\Delta_{\frac{7}{2}}^{7+}$	$\Delta_{\frac{9}{2}}^{9+}$	$\Delta_{\frac{11}{2}}^{11+}(2420)$
70	$\frac{1}{2}$	5		$N_{\frac{9}{2}}^{9-}$	$N_{\frac{11}{2}}^{11-}$	
	$\frac{3}{2}$	5	$N_{\frac{7}{2}}^{7-}$	$N_{\frac{9}{2}}^{9-}$	$N_{\frac{11}{2}}^{11-}(2600)$	$N_{\frac{13}{2}}^{13-}$

Table 1: $SU(6)$ multiplet structure for the known N and Δ baryon resonances including internal spin S and orbital angular momentum L quantum numbers. Radial excitations are non included in the table.

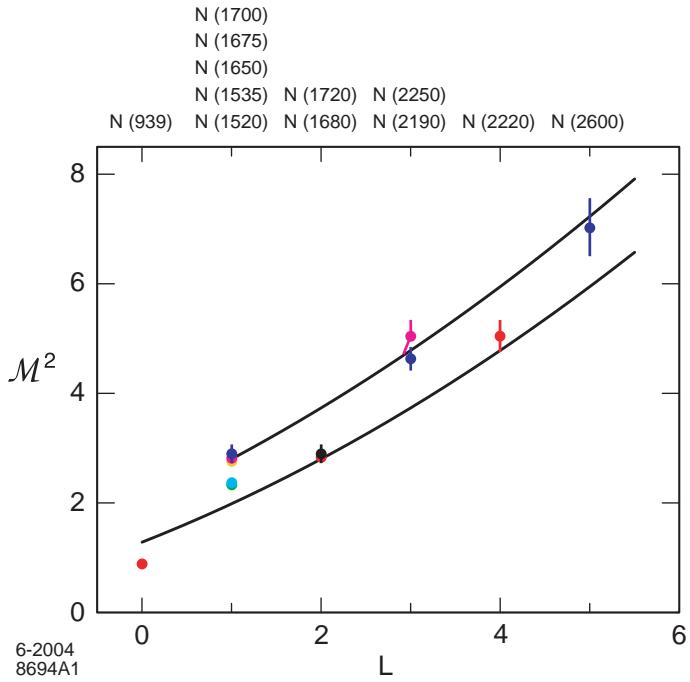


Figure 1: Nucleon orbital spectrum for a value of $\Lambda_{QCD} = 0.22$ GeV. The lower curve corresponds to nucleon states dual to spin- $\frac{1}{2}$ string modes in the bulk. The upper curve corresponds to nucleon states dual to string- $\frac{3}{2}$ modes.

actual parameter aside from the choice of the boundary conditions. The results for each trajectory show a clustering of states with the same orbital L , consistent with strongly suppressed spin-orbit forces; this is a severe problem for QCD models using one-gluon exchange. The results also indicate a parity degeneracy between states in the parallel trajectories shown in Fig. 1, as seen by displacing the upper curve by one unit of L to the right. Nucleon states with $S = \frac{3}{2}$ and Δ resonances fall on the same trajectory [26].

A string wavefunction with a node in the holographic coordinate z should correspond to a radial baryonic resonance with a node in the interquark separation, such as the positive parity Roper state $N_{\frac{1}{2}}^{1+}(1440)$. The first radial AdS eigenvalue corresponds to a mass around 1.85 GeV which would rather agree with the higher $N_{\frac{1}{2}}^{1+}(1710)$ state. Recent lattice results [30] are consistent with the $N_{\frac{1}{2}}^{1+}(1440)$ as a radial excited state. However, a possible interpretation as a $N_{\frac{1}{2}}^{1+}(1710)$ cannot be ruled out by present extrapolations. Lattice simulations at lower quark mass should be undertaken to have a definite answer.

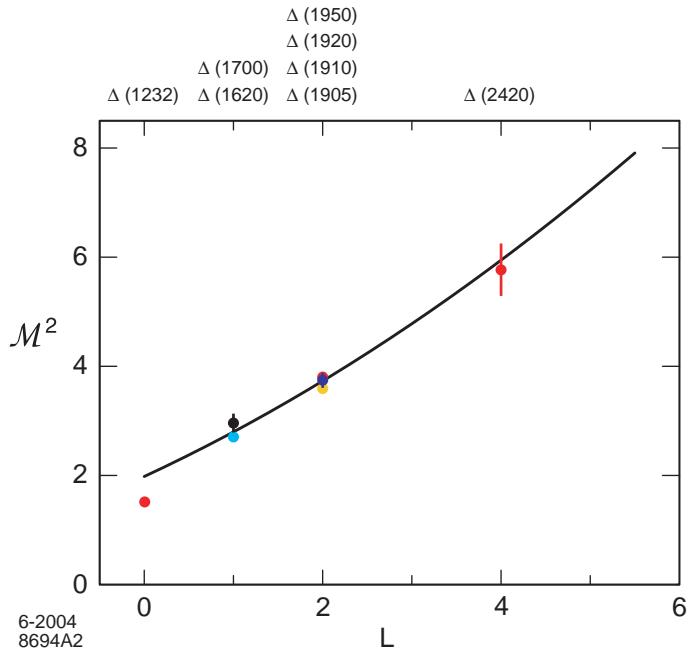


Figure 2: Delta orbital spectrum for $\Lambda_{QCD} = 0.22$ GeV. The Delta states dual to spin- $\frac{1}{2}$ and spin- $\frac{3}{2}$ string modes in the bulk lie on the same trajectory.

4 Concluding Remarks

We have described aspects of the $\mathcal{N} = 4$ SYM orbital baryon spectrum introducing an open string sector and a confining background by effectively cutting AdS space in the far infrared AdS region at $r_o = \Lambda_{QCD} R^2$. Since only one parameter, the QCD scale Λ_{QCD} , is used, the agreement of the model with the pattern of the physical light baryon spectrum is remarkable. This agreement possibly reflects the fact that our analysis is based on a conformal template, which is a good initial approximation to QCD [22]. We have chosen a special representation to construct a three quark baryon, and the results are effectively independent of N_C . This is consistent with results from lattice gauge theory for glueballs [31] where very little dependence on N_C at small lattice spacing is found for $N_C > 3$. The gauge/string correspondence presented here appears as a powerful organizing principle to classify and compute the mass eigenvalues of baryon resonances. A better understanding of nonconformal aspects of the metric and the nature of quantum fluctuations about the AdS geometry is required. Our results suggest that fundamental features of the hadron spectrum and QCD can be understood in terms of the nature of a higher dimensional dual gravity theory. Further discussion of some of these issues including a computation of the meson spectrum will be given elsewhere.

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Added Note

After completion of this work we have learned of several studies related to the present discussion: (1) A mapping of meson operators to fluctuations of the supergravity background has been conjectured in the framework of [20] in the low energy limit of closed strings. See: J. Erdmenger and I. Kirsh, arXiv:hep-th/0408113. (2) According to a recent proposal the lowest trajectory of Fig. 1 correspond to “good” diquarks, and the upper to “bad” diquarks. All the states in Fig. 2 correspond to “bad” diquarks. See: F. Wilczek, arXiv:hep-ph/0409168. (3) A recent lattice study of the Roper Resonance points out the necessity to carry further simulations at lower quark masses. See: D. Guadagnoli, M. Papinutto and S. Simula, arXiv:hep-lat/0409011.

References

- [1] J. M. Maldacena, *Adv. Theor. Math. Phys.* **2**, 231 (1998) [*Int. J. Theor. Phys.* **38**, 1113 (1999)] [arXiv:hep-th/9711200].
- [2] S. S. Gubser, I. R. Klebanov and A. M. Polyakov, *Phys. Lett. B* **428**, 105 (1998) [arXiv:hep-th/9802109].
- [3] E. Witten, *Adv. Theor. Math. Phys.* **2**, 253 (1998) [arXiv:hep-th/9802150].
- [4] J. Polchinski and M. J. Strassler, *Phys. Rev. Lett.* **88**, 031601 (2002) [arXiv:hep-th/0109174].
- [5] S. J. Brodsky and G. R. Farrar, *Phys. Rev. Lett.* **31**, 1153 (1973); *Phys. Rev. D* **11**, 1309 (1975); V. A. Matveev, R. M. Muradian and A. N. Tavkhelidze, *Lett. Nuovo Cim.* **7**, 719 (1973).
- [6] J. Polchinski and L. Susskind, arXiv:hep-th/0112204.
- [7] J. Polchinski and M. J. Strassler, *JHEP* **0305**, 012 (2003) [arXiv:hep-th/0209211].
- [8] H. Boschi-Filho and N. R. Braga, *Phys. Lett. B* **560**, 232 (2003) [arXiv:hep-th/0207071]; arXiv:hep-th/0312231.

- [9] R. C. Brower and C. I. Tan, Nucl. Phys. B **662**, 393 (2003) [arXiv:hep-th/0207144]; Nucl. Phys. Proc. Suppl. **119**, 938 (2003).
- [10] O. Andreev, Phys. Rev. D **67**, 046001 (2003) [arXiv:hep-th/0209256]; Phys. Rev. D **70**, 027901 (2004) [arXiv:hep-th/0402017].
- [11] For a review see, O. Aharony, S. S. Gubser, J. M. Maldacena, H. Ooguri and Y. Oz, Phys. Rept. **323**, 183 (2000) [arXiv:hep-th/9905111].
- [12] S. J. Brodsky and G. F. de Téramond, Phys. Lett. B **582**, 211 (2004) [arXiv:hep-th/0310227].
- [13] J. F. Gunion, S. J. Brodsky and R. Blankenbecler, Phys. Lett. B **39**, 649 (1972); Phys. Rev. **D8**, 287 (1973).
- [14] X. d. Ji, J. P. Ma and F. Yuan, Phys. Rev. Lett. **90**, 241601 (2003) [arXiv:hep-ph/0301141].
- [15] F. Bigazzi, A. L. Cotrone, M. Petrini and A. Zaffaroni, Riv. Nuovo Cim. **25N12**, 1 (2002) [arXiv:hep-th/0303191]; J. Russo, these proceedings.
- [16] W. Muck and K. S. Viswanathan, Phys. Rev. D **58**, 106006 (1998) [arXiv:hep-th/9805145].
- [17] I. R. Klebanov and E. Witten, Nucl. Phys. B **556**, 89 (1999) [arXiv:hep-th/9905104].
- [18] R. Camporesi and A. Higuchi, arXiv:gr-qc/9505009.
- [19] D. J. Gross and H. Ooguri, Phys. Rev. D **58**, 106002 (1998) [arXiv:hep-th/9805129]; E. Witten, JHEP **9807**, 006 (1998) [arXiv:hep-th/9805112].
- [20] A. Karch and E. Katz, JHEP **0206**, 043 (2002) [arXiv:hep-th/0205236]; B. A. Burrington, J. T. Liu, L. A. Pando Zayas and D. Vaman, arXiv:hep-th/0406207.
- [21] M. Kruczenski, D. Mateos, R. C. Myers and D. J. Winters, JHEP **0307**, 049 (2003) [arXiv:hep-th/0304032]; T. Sakai and J. Sonnenschein, JHEP **0309**, 047 (2003) [arXiv:hep-th/0305049]; D. Mateos, these proceedings.
- [22] S. J. Brodsky, arXiv:hep-ph/0408069.

- [23] S. S. Gubser, I. R. Klebanov and A. M. Polyakov, Nucl. Phys. B **636**, 99 (2002) [arXiv:hep-th/0204051].
- [24] A. Volovich, JHEP **9809**, 022 (1998) [arXiv:hep-th/9809009]; P. Matlock and K. S. Viswanathan, Phys. Rev. D **61**, 026002 (2000) [arXiv:hep-th/9906077].
- [25] S. Capstick and W. Roberts, Prog. Part. Nucl. Phys. **45**, S241 (2000) [arXiv:nucl-th/0008028].
- [26] E. Klempt, arXiv:hep-ph/0404270.
- [27] See for example, J. L. Goity, C. Schat and N. N. Scoccola, Phys. Lett. B **564**, 83 (2003) [arXiv:hep-ph/0304167]; N. Matagne, these proceedings.
- [28] LHP Collaboration; D. Richards, these proceedings.
- [29] S. Eidelman *et al.* [Particle Data Group Collaboration], Phys. Lett. B **592**, 1 (2004).
- [30] S. J. Dong, T. Draper, I. Horvath, F. X. Lee, K. F. Liu, N. Mathur and J. B. Zhang, arXiv:hep-ph/0306199; S. Sasaki, Prog. Theor. Phys. Suppl. **151**, 143 (2003) [arXiv:nucl-th/0305014]; T. Burch, C. Gattringer, L. Y. Glozman, R. Kleindl, C. B. Lang and A. Schaefer [Bern-Graz-Regensburg Collaboration], arXiv:hep-lat/0405006.
- [31] M. Teper, Nucl. Phys. Proc. Suppl. **109A**, 134 (2002) [arXiv:hep-lat/0112019]; M. Teper, these proceedings.